

Baryogenesis and the New Cosmology

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Abstract

In this talk I begin with a brief review of the status of approaches to understanding the origin of the baryon asymmetry of the universe (BAU). I then describe a recent model unifying three seemingly-distinct problems facing particle cosmology: the origin of inflation, the generation of the BAU and the nature of dark energy.

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I. INTRODUCTION

The audience for this talk was extremely diverse, ranging from beginning graduate students, through experts in subfields of physics somewhat distinct from the subject matter of my talk, to baryogenesis experts. My strategy therefore was to present an overview of the main issues facing baryogenesis, and only to focus on the more technical aspects of a particular model towards the end of the talk. Much of the review portion of these proceedings draws heavily on my article with Riotto [1] and the latter parts summarize my recent paper with De Felice and Nasri [2].

The problem of the baryon asymmetry of the universe is a classic problem of particle cosmology. Particle physics has taught us that matter and antimatter behave essentially identically, and indeed the interactions between matter and antimatter are the focus of successful terrestrial experiments. On the other hand, cosmology teaches us that the early universe was an extremely hot, and hence energetic, environment in which one would expect equal numbers of baryons and antibaryons to be copiously produced. This early state of the universe stands in stark contrast to what we observe in the universe today. Direct observation shows that the universe around us contains no appreciable primordial antimatter. In addition, the stunning success of big bang nucleosynthesis rests on the requirement that, defining $n_{b(\bar{b})}$ to be the number density of (anti)-baryons and s to be the entropy density,

$$2.6 \times 10^{-10} < \eta \equiv \frac{n_b - n_{\bar{b}}}{s} < 6.2 \times 10^{-10} , \quad (1)$$

(see, for example, [3]). Very recently this number has been independently determined to be $\eta = 6.1 \times 10^{-10} {}^{+0.3 \times 10^{-10}}_{-0.2 \times 10^{-10}}$ from precise measurements of the relative heights of the first two microwave background (CMB) acoustic peaks by the WMAP satellite [4]. Thus the natural question arises; as the universe cooled from early times to today, what processes, both particle physics and cosmological, were responsible for the generation of this very specific baryon asymmetry?

The outline of this talk is as follows. In the next section I survey several mechanisms by which the BAU might be generated. In particular I focus on electroweak baryogenesis and discuss why it is, to my mind, a particularly attractive mechanism. In section III I will then describe a model which attempts to unify three seemingly-disparate problems of modern cosmology; the origin of inflation, the generation of the BAU and the nature of dark energy.

II. OUR FAVORITE MODELS

If we're going to use a particle physics model to generate the baryon asymmetry of the universe (BAU), what properties must the theory possess? This question was first addressed by Sakharov [5] in 1967, resulting in the following criteria

- Violation of the baryon number (B) symmetry.
- Violation of the discrete symmetries C (charge conjugation) and CP (the composition of parity and C)
- A departure from thermal equilibrium.

Of course, the first of these is obvious - no B violation, no baryon production. To understand the second condition, note that, roughly speaking, if C and CP are conserved, the rate for any process which generates baryons is equal to that for the conjugate process, which produces antibaryons, so no net excess is generated on average. Finally, in thermal equilibrium the number density of a particle species is determined purely by its energy, and since the masses of particle and antiparticle are equal by the CPT theorem, the number density of baryons equals that of antibaryons. We will want to recall this connection between the equilibrium condition and the CPT theorem in the last part of this talk.

Grand Unified baryogenesis (for a review see [6]) was the first implementation of Sakharov's baryon number generation ideas. There are at least three reasons why this is an attractive mechanism. First, baryon number violation is an unavoidable consequence of Grand Unification. If one wishes to have quarks and leptons transforming in the same representation of a single group, then there necessarily exist processes in which quarks convert to leptons and vice-versa. Hence, baryon number violation is manifest in these theories. Second, it is a simple and natural task to incorporate sufficient amounts of CP-violation into GUTs. Among other possibilities, there exist many possible mixing phases that make this possible. Finally, at such early epochs of cosmic evolution, the relevant processes are naturally out of equilibrium since the relevant timescales are slow compared to the expansion rate of the universe. The simplicity with which the Sakharov criteria are achieved in GUTs makes this mechanism very attractive.

While GUT baryogenesis is attractive, it is not likely that the physics involved will be directly testable in the foreseeable future. While we may gain indirect evidence of grand unification with a particular gauge group, direct confirmation in colliders seems unrealistic. A second problem

with GUT scenarios is the issue of erasure of the asymmetry. As we will see, the electroweak theory ensures that there are copious baryon number violating processes between the GUT and electroweak scales. These events violate $B + L$ but conserve $B - L$. Thus, unless a GUT mechanism generates an excess $B - L$, any baryonic asymmetry produced will be equilibrated to zero by anomalous electroweak interactions. While this does not invalidate GUT scenarios, it is a constraint. For example, $SU(5)$ will not be suitable for baryogenesis for this reason, while $SO(10)$ may be. This makes the idea of baryogenesis through leptogenesis particularly attractive.

Affleck-Dine baryogenesis [7] is a particularly attractive scenario, and much progress has been made in understanding how this mechanism works. As was the case for electroweak baryogenesis, this scenario has found its most promising implementations [8, 9] in supersymmetric models, in which the necessary flat directions are abundant. Particularly attractive is the fact that these moduli, carrying the correct quantum numbers, are present even in the MSSM.

The challenges faced by Affleck-Dine models are combinations of those faced by the GUT and electroweak ideas. In particular, it is necessary that $B - L$ be violated along the relevant directions (except perhaps in the Q-ball implementations [10]) and that there exist new physics at scales above the electroweak. If supersymmetry is not found, then it is hard to imagine how the appropriate flat directions can exist in the low energy models.

In recent years, perhaps the most widely studied scenario for generating the baryon number of the universe has been electroweak baryogenesis. The popularity of this idea is tightly bound to its testability. The physics involved is all testable in principle at realistic colliders. Furthermore, the small extensions of the model involved to make baryogenesis successful can be found in supersymmetry, which is an independently attractive idea, although electroweak baryogenesis does not depend on supersymmetry. I will now describe this briefly.

A. Baryon Number Violation in the EW Theory

In the standard electroweak theory baryon number is an exact global symmetry. However, baryon number is violated at the quantum level through nonperturbative processes. These effects are closely related to the nontrivial vacuum structure of the electroweak theory.

At zero temperature, baryon number violating events are exponentially suppressed. This is because there exists a potential barrier between vacua and anomalous processes are thus tun-

neling events. The relevant barrier height is set by the point of least energy on the barrier. This point is known as the *sphaleron* [11], and has energy $E_{sph} \sim 10$ TeV. However, at temperatures above or comparable to the critical temperature of the electroweak phase transition, vacuum transitions over the sphaleron may occur frequently due to thermal activation [12].

B. C and CP Violation in the EW Theory

Fermions in the electroweak theory are chirally coupled to the gauge fields. In terms of the discrete symmetries of the theory, these chiral couplings result in the electroweak theory being maximally C-violating. However, the issue of CP-violation is more complex.

CP is known not to be an exact symmetry of the weak interactions, and is observed experimentally in the neutral Kaon system through K_0 , \bar{K}_0 mixing. Although at present there is no completely satisfactory theoretical explanation of this, CP violation is a natural feature of the standard electroweak model. The Kobayashi-Maskawa (KM) quark mass mixing matrix contains a single independent phase, a nonzero value for which signals CP violation. While this is encouraging for baryogenesis, it turns out that this particular source of CP violation is not strong enough. The relevant effects are parametrized by a dimensionless constant which is no larger than 10^{-20} . This appears to be much too small to account for the observed BAU and so it is usual to turn to extensions of the minimal theory. In particular the minimal supersymmetric standard model (MSSM).

C. The Electroweak Phase Transition

The question of the order of the electroweak phase transition is central to electroweak baryogenesis. Since the equilibrium description of particle phenomena is extremely accurate at electroweak temperatures, baryogenesis cannot occur at such low scales without the aid of phase transitions.

For a continuous transition, the associated departure from equilibrium is insufficient to lead to relevant baryon number production. The order parameter for the electroweak phase transition is φ , the modulus of the Higgs field. For a first order transition the extremum at $\varphi = 0$ becomes separated from a second local minimum by an energy barrier. At the critical temperature $T = T_c$ both phases are equally favored energetically and at later times the minimum at $\varphi \neq 0$ becomes the global minimum of the theory. Around T_c quantum tunneling

occurs and nucleation of bubbles of the true vacuum in the sea of false begins. At a particular temperature below T_c , bubbles just large enough to grow nucleate. These are termed *critical* bubbles, and they expand, eventually filling all of space and completing the transition. As the bubble walls pass each point in space, the order parameter changes rapidly, as do the other fields and this leads to a significant departure from thermal equilibrium. Thus, if the phase transition is strongly enough first order it is possible to satisfy the third Sakharov criterion in this way.

There is a further criterion to be satisfied. As the wall passes a point in space, the Higgs fields evolve rapidly and the Higgs VEV changes from $\langle\phi\rangle = 0$ in the unbroken phase to

$$\langle\phi\rangle = v(T_c) \tag{2}$$

in the broken phase. Here, $v(T)$ is the value of the order parameter at the symmetry breaking global minimum of the finite temperature effective potential. Now, CP violation and the departure from equilibrium occur while the Higgs field is changing. Afterwards, the point is in the true vacuum, baryogenesis has ended, and baryon number violation is exponentially suppressed. Since baryogenesis is now over, it is imperative that baryon number violation be negligible at this temperature in the broken phase, otherwise any baryonic excess generated will be equilibrated to zero. Such an effect is known as *washout* of the asymmetry and the criterion for this not to happen may be written as

$$\frac{v(T_c)}{T_c} \geq 1 . \tag{3}$$

Although there are a number of nontrivial steps that lead to this simple criterion, (3) is traditionally used to ensure that the baryon asymmetry survives after the wall has passed. It is necessary that this criterion be satisfied for any electroweak baryogenesis scenario to be successful.

In the minimal standard model, in which experiments at LEP now constrain the Higgs mass to be $m_H > 114.4$ GeV, it is clear from numerical simulations that (3) is not satisfied. This is therefore a second reason to turn to extensions of the minimal model.

In the MSSM there are two Higgs fields, Φ_1 and Φ_2 . At one loop, a CP-violating interaction between these fields is induced through supersymmetry breaking. Alternatively, there also exists extra CP-violation through CKM-like effects in the chargino mixing matrix. Thus, there seems to be sufficient CP violation for baryogenesis to succeed.

Now, the two Higgs fields combine to give one lightest scalar Higgs h . In addition, there are also light *stops* \tilde{t} (the superpartners of the top quark) in the theory. These light scalar particles

can lead to a strongly first order phase transition if the scalars have masses in the correct region of parameter space. A detailed two loop calculation [13] and lattice results indicate that the allowed region is given by

$$m_h \leq 120\text{GeV} \quad (4)$$

$$m_{\tilde{t}} \leq m_t, \quad (5)$$

for $\tan\beta \equiv \langle\Phi_2\rangle/\langle\Phi_1\rangle > 5$. In the next few years, experiments at the Tevatron and the LHC should probe this range of Higgs masses and we should know if the MSSM is a good candidate for electroweak baryogenesis.

Thus, the testability of electroweak scenarios also leads to tight constraints, and at present there exists only a small window of parameter space in extensions of the electroweak theory in which baryogenesis is viable.

If the Higgs is found, the second test will come from the search for the lightest stop at the Tevatron collider. If both particles are found, the last crucial test will come from CP-violating effects which may be observable in B physics. Moreover, the preferred parameter space leads to values of the branching ratio $\text{BR}(b \rightarrow s\gamma)$ different from the Standard Model case. Although the exact value of this branching ratio depends strongly on the value of the μ and A_t parameters, the typical difference with respect to the Standard Model prediction is of the order of the present experimental sensitivity and hence in principle testable in the near future. Indeed, for the typical spectrum considered here, due to the light charged Higgs, the branching ratio $\text{BR}(b \rightarrow s\gamma)$ is somewhat higher than in the SM case, unless negative values of $A_t\mu$ are present.

Having given an overview of baryogenesis and a description of a particularly attractive mechanism, I will turn, in the rest of my talk, to a mechanism I've been working on recently - *quintessential baryogenesis* [2].

III. QUINTESSENTIAL BARYOGENESIS

Rolling scalar fields are a mainstay of modern cosmology. This is perhaps best-illustrated by the inflationary paradigm [14, 15, 16], in which most implementations involve a scalar field rolling towards the minimum of its potential in such a way that the potential energy of the field is the dominant component of the energy density of the universe. There are, however, many other cosmological instances in which scalar fields are invoked. In fact, it is often felt that

cosmologists are prepared to invent a new scalar field every time a new piece of cosmological data comes along.

During the last few years a new consistent picture of the energy budget of the universe has emerged. Large scale structure studies show that matter (both luminous and dark) contributes a fraction of about 0.3 of the critical density, while the position of the first acoustic peak of the cosmic microwave background power spectrum indicates that the total energy density is consistent with criticality. The discrepancy between these two measurements may be reconciled by invoking a negative pressure component which is termed *dark energy* and leads to the acceleration of the universe. Current limits [17], obtained by combining results from cosmic microwave background experiments with large scale structure data, the Hubble parameter measurement from the Hubble Space Telescope and luminosity measurements of Type Ia supernovae, give $-1.62 < w < -0.74$ at the 95% confidence level.

It is of course possible that this mystery component is a cosmological constant Λ , for which $w_\Lambda = -1$. Alternatively, it has been suggested [18]-[20] that if the cosmological constant itself is zero, the dark energy component could be due to the dynamics of a rolling scalar field, in a form of late-universe inflation that has become known as *quintessence*.

It is natural to wonder whether the inflaton and the quintessence field might be one and the same [21], and, in fact, specific models for this have been proposed [21]-[24]. In this part of my talk I'll investigate how we may further limit the proliferation of rolling scalar fields required in modern cosmology by studying how a scalar field responsible both for inflation and for late-time acceleration of the universe might also be responsible for the generation of the baryon asymmetry of the universe. The relationship between early-time acceleration – inflation – and baryogenesis has been explored in some detail (for example see [7], [25]-[38]). Here I consider the opposite regime, that the quintessence field may be associated with the generation of the baryon asymmetry [2] (for interesting related studies see [39, 40, 41]).

A. Quintessential Inflation

I'll use the flat Friedmann, Robertson-Walker (FRW) ansatz for the metric

$$ds^2 = -dt^2 + a(t)^2 \left[dr^2 + r^2 d\Omega_2^2 \right] . \quad (6)$$

Here the energy density ρ and pressure p for a real homogeneous scalar field ϕ are given by

$$\rho_\phi = \frac{1}{2} \dot{\phi}^2 + V(\phi) , \quad (7)$$

$$p_\phi = \frac{1}{2}\dot{\phi}^2 - V(\phi) , \quad (8)$$

respectively, with $V(\phi)$ the potential, and where we have defined the Planck mass by $G \equiv M_{\text{p}}^{-2}$. The scalar field itself obeys

$$\ddot{\phi} + 3 \left(\frac{\dot{a}}{a} \right) \dot{\phi} + \frac{dV(\phi)}{d\phi} = 0 , \quad (9)$$

with a prime denoting a derivative with respect to ϕ .

Now, to explain the current data indicating an accelerating universe, it is necessary to have the dominant type of matter at late times be such that $\ddot{a} > 0$. If this matter is to be ϕ , then this implies $\rho_\phi + 3p_\phi < 0$. Writing $p_\phi \equiv w_\phi \rho_\phi$ this translates into an equation of state parameter that obeys $w_\phi < -\frac{1}{3}$

In the Peebles and Vilenkin model [22] one uses the potential

$$V(\phi) = \begin{cases} \lambda(\phi^4 + M^4) & , \quad \phi \in (-\infty, 0] \\ \frac{\lambda M^8}{\phi^4 + M^4} & , \quad \phi \in (0, \infty) , \end{cases} \quad (10)$$

We require inflation to occur at early times, and indeed, for this potential, the slow-roll conditions are satisfied for sufficiently large and negative ϕ .

Inflation ends when the slow-roll conditions are violated and the potential and kinetic energies of the inflaton are comparable with each other. In traditional inflationary models the inflaton then rapidly transfers its energy to other fields either through perturbative effects (reheating) or parametric resonance (preheating). Here, however, there is no such effect, and it is the kinetic energy of the field ϕ that is the dominant component of the energy density of the universe immediately after the end of inflation. Following Joyce [42] we term this behavior *kination*.

A successful cosmology requires the universe be radiation-dominated at the time of nucleosynthesis, since otherwise the precision predictions of that theory are no longer in agreement with observations. The lack of conventional reheating, the conversion of the potential energy of the inflaton to particle production, in quintessential inflation means that the requisite radiation must be produced another way. In fact, the radiation era in these models is due to the subtle behavior of quantum fields in changing geometries.

At the end of inflation, the FRW line element undergoes an abrupt change from that associated with cosmic expansion (exponential or power-law) to that associated with kination. Massless quantum fields in their vacua in the inflation era are no longer in vacuum in the kination era, corresponding to gravitational particle production. This effect is analogous to

Hawking radiation, and has been explored in detail [43]-[45] in the cosmological context of interest here.

For a significant time subsequent to this, cosmic evolution is much the same as in the standard cosmology, with a matter dominated epoch eventually succeeding the radiation era. Although the scalar field is not important during these times, the density fluctuations seeded by quantum fluctuations in ϕ during inflation lead to structure formation and temperature fluctuations in the cosmic microwave background radiation.

Finally, consider this extreme future of the universe, in which the scalar field becomes responsible for quintessence. It is clear that when quintessence begins, the energy density of the field ϕ once again becomes dominated by its potential energy density. It is a challenge similar to that for conventional quintessence to ensure that this epoch occurs at the present time and yields the correct ratio of matter to dark energy.

B. Generation of the Baryon Asymmetry

In order for the quintessence field ϕ to play a role in baryogenesis, we must consider how ϕ couples to other fields. In principle, the inflaton and quintessence field may lie in any sector of the theory, the phenomenologically safest of which would be one in which there are only gravitational strength couplings to other particles. Here I adopt a conservative approach and assume that ϕ couples to standard model fields with couplings specified by a dimensionless constant and an energy scale which we shall leave as a free parameter for the moment and later constrain by observations and the condition that our model produce a sufficient baryon asymmetry.

Consider terms in the effective Lagrangian density of the form

$$\mathcal{L}_{\text{eff}} = \frac{\lambda'}{M} \partial_\mu \phi J^\mu, \quad (11)$$

where λ' is a coupling constant, $M < M_p$ is the scale of the cutoff in the effective theory and J^μ is the current corresponding to some continuous global symmetry such as baryon number or baryon number minus lepton number. Further, let us assume that ϕ is homogeneous. We then obtain

$$\mathcal{L}_{\text{eff}} = \frac{\lambda'}{M} \dot{\phi} \Delta n \equiv \mu(t) \Delta n, \quad (12)$$

where $n = J^0$ is the number density corresponding to the global symmetry and we have defined an effective time-dependent “chemical potential” $\mu(t) \equiv \lambda' \dot{\phi} / M$.

Recall that we need to satisfy the Sakharov criteria in order to generate a baryon asymmetry (for reviews see [46]-[48], [1]). The first of these requires baryon number B to be violated. At this stage, to maintain generality, we shall leave the mechanism of baryon number violation unspecified. Further, the standard model is maximally C -violating due to its chiral structure, and the coupling (11) is CP -odd. In this sense, no *explicit* CP -violation is required in this model. The third Sakharov criterion requires a departure from thermal equilibrium if CPT is a manifest symmetry. However, the crucial point about baryogenesis in the presence of the rolling scalar field ϕ is that CPT is broken spontaneously by the explicit value taken by $\langle\dot{\phi}\rangle \neq 0$. Thus, the particular model of baryogenesis that is important here is *spontaneous baryogenesis* [49], which is effective even in thermal equilibrium. Following [2], I will refer to this model, in which the rolling scalar responsible for both inflation and dark energy also provides a source for spontaneous baryogenesis as *quintessential baryogenesis*.

To understand how spontaneous (and hence quintessential) baryogenesis works, note that in thermal equilibrium we have

$$\Delta n(T; \xi) = \int \frac{d^3\mathbf{p}}{(2\pi)^3} [f(E, \mu) - f(E, -\mu)] , \quad (13)$$

where $\xi \equiv \mu/T$ is a parameter and $f(E, \mu)$ is the phase-space distribution of the particles of the current J^μ , which may be Fermi-Dirac or Bose-Einstein. Thus, for $\xi < 1$

$$\Delta n(T; \mu) \simeq \frac{gT^3}{6}\xi + \mathcal{O}(\xi^2) , \quad (14)$$

where g is the number of degrees of freedom of the field corresponding to n . Therefore,

$$\Delta n(T; \mu) \simeq \frac{\lambda' g}{6M} T^2 \dot{\phi} . \quad (15)$$

Now recall that the entropy density is given by

$$s = \frac{2\pi}{45} g_* T^3 , \quad (16)$$

where g_* is the effective number of relativistic degrees of freedom in thermal equilibrium at temperature T . Whatever the mechanism of baryon number violation, there will exist a temperature T_F below which baryon number violating processes due to this mechanism become sufficiently rare that they freeze out. For $T < T_F$ these processes can no longer appreciably change the baryon number of the universe. Computing the freeze-out value of the baryon to entropy ratio we then obtain

$$\eta_F \equiv \eta(T_F) \equiv \frac{\Delta n}{s}(T_F) \simeq 0.38 \lambda' \left(\frac{g}{g_*} \right) \frac{\dot{\phi}(T_F)}{MT_F} . \quad (17)$$

How the baryon excess evolves after this point depends on the value of T_F and on the relevant current in equation (11). If $T_F \leq T_c^{\text{EW}} \sim 100$ GeV, the critical temperature of the electroweak phase transition, then all baryon number violation ceases at T_F and $\eta(T < T_F) = \eta_F$. However, if $T_F > T_c^{\text{EW}}$, then we must take into account the effects of anomalous electroweak processes at finite temperature. These can be involved in directly generating the baryon asymmetry (electroweak baryogenesis), in reprocessing an asymmetry in other quantum numbers into one in baryon number (for example in leptogenesis) or in diluting the asymmetry created by any baryogenesis mechanism which is effective above the electroweak scale and does not produce a $B - L$ asymmetry. It is important to realize that, in the context of quintessential inflation, the quantitative effects of these electroweak processes may differ substantially from those in the standard cosmology [42, 50] since in our case, the electroweak phase transition may occur during kination rather than radiation domination.

I have now provided quite a general description of quintessential baryogenesis, appropriate to the level of this talk. While this allows us to demonstrate the generic features of our model, we cannot calculate the magnitude of the actual baryon asymmetry generated without first specifying a mechanism of baryon number violation (and hence a value for T_F) and a value for the dimensionless combination $\lambda' M_{\text{p}}/M$. Here I'll provide just one concrete example.

C. Baryon Number Violation Through Non-renormalizable Operators

If there exists baryon number violating physics above the standard model, then this physics will manifest itself in non-renormalizable operators in the standard model. For the purposes of this section we will actually be interested in operators that violate the anomaly-free combination $B - L$. Consider the effective 4-fermion operator

$$\mathcal{L}_{B-L} = \frac{\tilde{g}}{M_X^2} \psi_1 \psi_2 \bar{\psi}_3 \bar{\psi}_4 , \quad (18)$$

where ψ_i denote standard model fermions. Here \tilde{g} is a dimensionless coupling, obtained after integrating out the $B - L$ violating effects of a particle of mass M_X . The rate of baryon number violating processes due to this operator is, as usual, defined by $\Gamma_{B-L}(T) = \langle \sigma(T) n(T) v \rangle$, where $\sigma(T)$ is the cross-section for $\psi_1 + \psi_2 \rightarrow \psi_3 + \psi_4$, $n(T)$ is the number density of ψ particles, v is the relative velocity and $\langle \dots \rangle$ denotes a thermal average. For temperatures $T < M_X$ we have $n(T) \sim T^3$, $\sigma(T) \sim \tilde{g}^2 T^2 / M_X^4$, and $v \sim 1$ which yields

$$\Gamma_{B-L}(T) \simeq \frac{\tilde{g}^2}{M_X^4} T^5 . \quad (19)$$

The high power of the temperature dependence in this rate results from the fact that (18) is an irrelevant operator in the electroweak theory and is crucial for the success of our mechanism. These interactions are in thermal equilibrium in the early universe, but because their rate drops off so quickly with the cosmic expansion they will drop out of equilibrium at the temperature T_F defined through

$$\Gamma_{B-L}(T_F) = H(T_F) . \quad (20)$$

As a definite example, let us take $\lambda' M_p/M \sim 8$, which is as large as is allowed by current constraints. In this case we need $M_X^{\text{PV}} \sim 10^{11}$ GeV, the intermediate scale that appears in some supersymmetric models.

D. Brief Comments on Constraints and Tests

The presence of an extremely light scalar field in the universe has the potential to lead to a number of observable consequences in the laboratory and in cosmology. In the case of quintessence these effects have been analyzed in some detail by Carroll [51]. Particularly strong constraints arise due to couplings of the form

$$\mathcal{L}_{\text{eff}}^{(1)} \equiv \beta_{F\tilde{F}} \frac{\phi}{M} F_{\mu\nu} \tilde{F}^{\mu\nu} , \quad (21)$$

where $F_{\mu\nu}$ and $\tilde{F}^{\mu\nu} \equiv \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$ are the electromagnetic field strength tensor and its dual respectively. If, as in quintessence models, the field ϕ is homogeneous and time varying, then it affects the dispersion relation for electromagnetic waves and leads to a rotation in the direction of polarized light from radio sources [52, 53]. To avoid such bounds, quintessence models usually struggle to have such couplings be as small as possible. In our model however, we are making important use of the coupling (11). If the relevant current is that for baryon number J_B^μ , then it can be shown that Carroll's bound yields $\lambda' M_p/M < 8$. If the relevant current is J_{B-L}^μ then the above argument does not apply formally. However, we may still generally expect an analogous coupling to the term $F^{\mu\nu} \tilde{F}_{\mu\nu}$ of a similar order.

We have already demonstrated that, for an appropriate scale M_X , successful quintessential baryogenesis takes place for $\lambda' < 8$, and so it is possible to generate the observed BAU and to evade existing constraints. If quintessential baryogenesis is correct, then it may be that the relevant coupling λ' lies just below the existing observational bounds and that future studies of the rotation of polarized light from distant galaxies will reveal the presence of such a term.

IV. COMMENTS AND CONCLUSIONS

In the first part of this talk I presented a current view of the status of baryogenesis models, focusing on the best-known ones.

In this second part of my talk I extended the approach of Spokoiny [21] and of Peebles and Vilenkin [22] in exploring the extent to which the dynamics of a single scalar field can be responsible for not only inflation and dark energy domination, but also the baryon asymmetry of the universe. This model is called quintessential baryogenesis [2], and is interesting, of course, since the best fit cosmology to all current observational data is one in which the universe undergoes two separate epochs of acceleration.

We have left a number of questions unanswered and will return to them in future work. Perhaps the most pressing issue is one that plagues rolling scalar models of dark energy in general, namely the question of technical naturalness of the potentials involved, and their stability to quantum corrections. However, this is a general issue for quintessence models, and is not specific to our baryogenesis mechanism.

One thing is certain: the existence of the BAU is clear evidence from cosmology of particle physics beyond the standard model. This makes the conundrum of the BAU a fascinating problem for cosmologists and high-energy physicists alike.

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